

750 GeV Diphoton Resonance from Singlets in an Exceptional Supersymmetric Standard Model

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Abstract

The 750-760 GeV diphoton resonance may be identified as one or two scalars and/or one or two pseudoscalars contained in the two singlet superfields $S_{1,2}$ arising from the three 27-dimensional representations of E_6 . The three 27s also contain three copies of colour-triplet charge $\mp 1/3$ vector-like fermions D, \bar{D} and two copies of charged inert Higgsinos \tilde{H}^+, \tilde{H}^- to which the singlets $S_{1,2}$ may couple. We propose a variant of the E_6 SSM where the third singlet S_3 breaks a gauged $U(1)_N$ above the TeV scale, predicting $Z'_N, D, \bar{D}, \tilde{H}^+, \tilde{H}^-$ at LHC Run 2, leaving the two lighter singlets $S_{1,2}$ with masses around 750 GeV. We calculate the branching ratios and cross-sections for the two scalar and two pseudoscalar states associated with the $S_{1,2}$ singlets, including possible degeneracies and maximal mixing, subject to the constraint that their couplings remain perturbative up to the unification scale.

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1 Introduction

Recently ATLAS and CMS experiments have reported an excess of diphoton events at an invariant mass around 750 GeV and 760 GeV from LHC Run-2 with pp collisions at the center of mass energy of 13 TeV [1, 2]. The local significance of the excess ATLAS events is 3.9σ while that of the excess CMS events is 2.6σ , corresponding to the respective cross sections $\sigma(pp \rightarrow \gamma\gamma) = 10.6$ fb and $\sigma(pp \rightarrow \gamma\gamma) = 6.3$ fb. ATLAS favours a broad width of $\Gamma \sim 45$ GeV, while CMS, although not excluding a broad resonance, actually prefers a narrow width. The diphoton excesses observed by ATLAS and CMS at this mass scale may be partially understood by the factor of 5 gain in cross-section due to gluon production. However there is no evidence for any coupling of the resonance into anything except gluons and photons (no final states such as $f\bar{f}$, VV (f being a fermion and V being W, Z) since no missing E_T or jets have been observed.

This may be the first indication of new physics at the TeV scale. It could even be the tip of an iceberg of many future discoveries. Several interpretations have been suggested based on extensions of the Standard Model spectrum [3]-[157]. Many of these papers suggest a spinless singlet coupled to vector-like fermions [3, 9, 10, 12, 14, 21, 22, 34, 37, 55, 61, 63, 83, 84, 98, 104, 107, 109, 157, 125, 121, 131]. Indeed, the observed resonance could be interpreted as a Standard Model scalar or pseudoscalar singlet state X with mass $m_X \sim 750 - 760$ GeV. Moreover, because it decays into two photons, its spin is consistent with $s = 0$. The process of generating the two photons can take place by the gluon-gluon fusion mechanism according to the process $gg \rightarrow X \rightarrow \gamma\gamma$ hence it requires production and decay of the particle X . In a renormalisable theory this interaction can be realised assuming vector-like fermions at the TeV scale, which carry electric charge and colour. Such vector-like pairs have not been observed at LHC, hence the mass of the fermion pair should be around or above the TeV scale. For example, in F-theory models based on E_6 , low energy singlets coupling to extra vector-like matter is predicted and may be responsible for the 750 GeV diphoton resonance [157]. Such models motivate the phenomenological study of E_6 as being the origin of the new physics.

An example of a model with singlets and extra vector-like matter is the Exceptional Supersymmetric (SUSY) Standard Model (E_6 SSM) [158, 159], where the spectrum of the MSSM is extended to fill out three complete 27-dimensional representations of the gauge group E_6 which is broken at the unification scale down to the SM gauge group plus an additional gauged $U(1)_N$ symmetry at low energies under which right-handed neutrinos are neutral, allowing them to get large masses. The three 27_i -plet families (labelled by $i = 1, 2, 3$) contain the usual quarks and leptons plus the following extra states: SM-singlet fields, S_i ; up- and down-type Higgs doublets, H_{ui} and H_{di} ; and charged $\pm 1/3$ coloured, exotics D_i and \bar{D}_i . The extra matter ensures anomaly cancellation, however the model also contains two extra $SU(2)$ doublets, H' and \bar{H}' , which are required for gauge coupling unification [160]. To evade rapid proton decay a \mathbb{Z}_2 symmetry, either \mathbb{Z}_2^{qq} or \mathbb{Z}_2^{lq} , is introduced and to evade large flavour changing neutral currents an approximate \mathbb{Z}_2^H symmetry is introduced where only the third family of Higgs doublets H_{u3} and H_{d3} and singlets S_3 are even under it and hence couple to fermions and get vacuum expectation values (VEVs). In particular, the third family singlet S_3 gets a VEV, $\langle S_3 \rangle = s/\sqrt{2}$, which is responsible for the effective μ term, inert Higgsino and D-fermion and Z'_N masses, while the first and second families of Higgs doublets and SM-singlets do not get VEVs and are called “inert”. Further aspects of the theory and phenomenology of this SUSY extension of the SM have been extensively studied in [161, 162, 163, 164, 165, 166, 167, 168, 169].

In this paper we take all three singlets to be even under the approximate \mathbb{Z}_2^H , which allows them all to couple to \hat{H}_{ui} and \hat{H}_{di} as well as \hat{D}_i and $\hat{\bar{D}}_i$. We shall assume that the third singlet S_3 has appreciable couplings to the three families of H_{ui} , H_{di} and D_i , \bar{D}_i , so that its large VEV generates effective mass terms for all these states, as well a Z'_N , above the TeV scale but possibly

within the reach of LHC Run 2. However the first and second singlets $S_{1,2}$ may have relatively small couplings to the third pair of Higgs doublets H_{u3} and H_{d3} , which are the only Higgs doublets to acquire VEVs. In addition, we shall suppose that the value of the third singlet S_3 VEV s is above the TeV scale, while the other singlets $S_{1,2}$ at most develop small VEVs. This is different from the modified E_6 SSM in [125], where two of the singlets were assigned even under the approximate \mathbb{Z}_2^H and both were allowed to develop VEVs and couple to all three families of H_{ui}, H_{di} and D_i, \bar{D}_i . In the version of the E_6 SSM here, we suppose that, after the third singlet S_3 with large s VEV is integrated out, only the first and second singlets $S_{1,2}$ appear in the low energy effective theory and provide candidates for the 750-760 GeV resonance which may be identified as one or two scalars and/or one or two pseudoscalars contained in $S_{1,2}$. The assumed smallness of the coupling of $S_{1,2}$ to H_{u3} and H_{d3} means that the observed resonance will not easily decay into pairs of top quarks or W bosons.

Many of the features of the considered model would be common to other SUSY E_6 models where the low energy spectrum consists of complete 27-plets. The present model is a variant of the E_6 SSM and like that model is distinguished by the choice of surviving gauged $U(1)_N$ under which right-handed neutrinos have zero charge and may acquire large Majorana masses, corresponding to a high scale seesaw mechanism. For earlier literature on other SUSY E_6 models based on different surviving gauged $U(1)$ symmetries under which right-handed neutrinos are charged see [158].

The layout of the remainder of the paper is as follows. In section 2 we discuss the variant of the E_6 SSM that we shall study, and discuss the renormalisation group equations which constrain the Yukawa couplings to be perturbative up to the unification scale. In section 3 we apply this model to the 750 GeV diphoton resonance, calculating the branching ratios and cross-sections for the two scalar and two pseudoscalar states associated with the $S_{1,2}$ singlets, including possible degeneracies and mixing. Section 4 concludes the paper.

2 A variant of the E_6 SSM

We first recall that the E_6 SSM [158, 159] may be derived from an E_6 GUT group broken via the following symmetry breaking chain:

$$\begin{aligned} E_6 &\rightarrow SO(10) \otimes U(1)_\psi \\ &\rightarrow SU(5) \otimes U(1)_\chi \otimes U(1)_\psi \\ &\rightarrow SU(3) \otimes SU(2) \otimes U(1)_Y \times U(1)_\chi \otimes U(1)_\psi \\ &\rightarrow SU(3) \otimes SU(2) \otimes U(1)_Y \otimes U(1)_N. \end{aligned} \tag{1}$$

We assume that the above symmetry breaking chain occurs at a single GUT scale M_X in one step, due to some unspecified symmetry breaking sector,

$$E_6 \rightarrow SU(3) \otimes SU(2) \otimes U(1)_Y \otimes U(1)_N, \tag{2}$$

where

$$U(1)_N = \cos(\vartheta)U(1)_\chi + \sin(\vartheta)U(1)_\psi \tag{3}$$

and $\tan(\vartheta) = \sqrt{15}$ such that the right-handed neutrinos that appear in the model are completely neutral and may get large intermediate scale masses. However the $U(1)_N$ gauge group remains unbroken down to the few TeV energy scale where its breaking results in an observable Z'_N . Three complete 27 representations of E_6 then also must survive down to this scale in order

to ensure anomaly cancellation. These 27_i decompose under the $SU(5) \otimes U(1)_N$ subgroup as follows:

$$27_i \rightarrow (10, 1)_i + (\bar{5}, 2)_i + (\bar{5}, -3)_i + (5, -2)_i + (1, 5)_i + (1, 0)_i, \quad (4)$$

where the $U(1)_N$ charges must be GUT normalised by a factor of $1/\sqrt{40}$. The first two terms contain the usual quarks and leptons, and the final term, which is a singlet under the entire low energy gauge group, contains the (CP conjugated) right-handed neutrinos N_i^c . The last-but-one term, which is charged only under $U(1)_N$, contains the SM-singlet fields S_i . The remaining terms $(\bar{5}, -3)_i$ and $(5, -2)_i$ contain three families of up- and down-type Higgs doublets, H_{ui} and H_{di} , and charged $\pm 1/3$ coloured exotics, D_i and \bar{D}_i . These are all superfields and are written with hats in the following.

The low energy gauge invariant superpotential can be written

$$W^{\text{E}_6\text{SSM}} = W_0 + W_{1,2}, \quad (5)$$

where $W_{0,1,2}$ are given by

$$W_0 = \lambda_{jki} \hat{H}_{dj} \hat{H}_{uk} \hat{S}_i + \kappa_{jki} \hat{D}_j \hat{D}_k \hat{S}_i + h_{ijk}^N \hat{N}_i^c \hat{H}_{uj} \hat{L}_{Lk} \\ + h_{ijk}^U \hat{H}_{ui} \hat{Q}_{Lj} \hat{u}_{Rk}^c + h_{ijk}^D \hat{H}_{di} \hat{Q}_{Lj} \hat{d}_{Rk}^c + h_{ijk}^E \hat{H}_{di} \hat{L}_{Lj} \hat{e}_{Rk}^c, \quad (6)$$

$$W_1 = g_{ijk}^Q \hat{D}_i \hat{Q}_{Lj} \hat{Q}_{Lk} + g_{ijk}^q \hat{D}_i \hat{d}_{Rj}^c \hat{u}_{Rk}^c, \quad (7)$$

$$W_2 = g_{ijk}^N \hat{N}_i^c \hat{D}_j \hat{d}_{Rk}^c + g_{ijk}^E \hat{D}_i \hat{u}_{Rj}^c \hat{e}_{Rk}^c + g_{ijk}^D \hat{D}_i \hat{Q}_{Lj} \hat{L}_{Lk}, \quad (8)$$

with $W_{1,2}$ referring to either W_1 or W_2 (but not both together which would result in excessive proton decay unless the associated Yukawa couplings were very small).

At the renormalisable level the gauge invariance ensures matter parity and hence LSP stability. All lepton and quark superfields are defined to be odd under matter parity \mathbb{Z}_2^M , while \hat{H}_{ui} , \hat{H}_{di} , \hat{D}_i , $\hat{\bar{D}}_i$, and \hat{S}_i are even. This means that the fermions associated with \hat{D}_i , $\hat{\bar{D}}_i$ are SUSY particles analogous to the Higgsinos, while their scalar components may be thought of as colour-triplet (and electroweak singlet) Higgses, making complete 5 and $\bar{5}$ representations without the usual doublet-triplet splitting.

In order for baryon and lepton number to also be conserved, preventing rapid proton decay mediated by \hat{D}_i , $\hat{\bar{D}}_i$, one imposes either \mathbb{Z}_2^{qq} or \mathbb{Z}_2^{lq} . Under \mathbb{Z}_2^{qq} , the lepton, including the RH neutrino, superfields are assumed to be odd, which forbids W_2 . Under \mathbb{Z}_2^{lq} , the lepton and the \hat{D}_i , $\hat{\bar{D}}_i$ superfields are assumed odd, which forbids W_1 . Baryon and lepton number are conserved at the renormalisable level, with the \hat{D}_i , $\hat{\bar{D}}_i$ interpreted as being either diquarks in the former case or leptoquarks in the latter case.

In the $E_6\text{SSM}$, a further approximate flavour symmetry \mathbb{Z}_2^H was also assumed. It is this approximate symmetry that distinguishes the third (by definition, “active”) generation of Higgs doublets and SM-singlets from the second and first (“inert”) generations. Under this approximate symmetry, all superfields are taken to be odd, apart from the active \hat{S}_3 , \hat{H}_{d3} , and \hat{H}_{u3} which are taken to be even. The inert fields then have small couplings to matter and do not radiatively acquire VEVs or lead to large flavour changing neutral currents. The active fields can have large couplings to matter and radiative electroweak symmetry breaking (EWSB) occurs with these fields. In particular the VEV $\langle S_3 \rangle = s/\sqrt{2}$ is responsible for breaking the $U(1)_N$ gauge group and generating the effective μ term and D-fermion masses. In particular we must have $s > 5$ TeV in order to satisfy $M_{Z'_N} > 2.5$ TeV, which is the current LHC Run 2 experimental limit [170].

We now propose a variant of the E₆SSM in which we allow all three singlets \hat{S}_i (as well as \hat{H}_{d3} and \hat{H}_{u3}) to be even under the \mathbb{Z}_2^H . This allows all three singlets \hat{S}_i to couple to \hat{H}_{ui} and \hat{H}_{di} as well as \hat{D}_i and $\hat{\bar{D}}_i$. If for simplicity we take the couplings in Eq. (6) to have the diagonal form, $\lambda_{jki} \propto \lambda_{ji}\delta_{jk}$ and $\kappa_{jki} \propto \kappa_{ji}\delta_{jk}$, then the \mathbb{Z}_2^H symmetry allows to reduce the structure of the Yukawa interactions in the superpotential to:

$$W^{\text{E}_6\text{SSM}} \simeq \lambda_{ji}\hat{H}_{dj}\hat{H}_{uj}\hat{S}_i + \kappa_{ji}\hat{D}_j\hat{\bar{D}}_j\hat{S}_i + W_{MSSM}(\mu = 0). \quad (9)$$

The superfield \hat{S}_3 is assumed to acquire a rather large VEV ($\langle S_3 \rangle = s/\sqrt{2}$) giving rise to the effective μ term, masses of exotic quarks and inert Higgsino states which are given by

$$\mu = \frac{\lambda_{33}s}{\sqrt{2}}, \quad \mu_{H_\alpha} = \frac{\lambda_{\alpha 3}s}{\sqrt{2}}, \quad \mu_{D_i} = \frac{\kappa_{i3}s}{\sqrt{2}},$$

In our analysis here we restrict our consideration to the case when exotic quarks and inert Higgsinos are sufficiently light compared to the VEV $s > 5$ TeV, but are heavier than half the mass of the 750 GeV resonance, so that they appear in loop diagrams for the singlet decays. It means that the Yukawa couplings of \hat{S}_3 to all exotic states should be quite small. Throughout this paper we are going to assume that some scalar components of the first and second singlets \hat{S}_α , with $\alpha = 1, 2$, can be identified with the resonances which give rise to the excess of diphoton events at an invariant mass around 750 GeV recently reported by the LHC experiments. ATLAS and CMS measurements indicate that the branching ratios of the decays of such resonances into SM fermions have to be sufficiently small. This implies that the mixing between the scalar components of \hat{S}_α and the neutral scalar components of the third pair of Higgs doublets H_u and H_d , which are the ones that give rise to the EWSB, should be strongly suppressed. In order to ensure the suppression of the corresponding mixing we impose the further requirement, namely that the SM singlets \hat{S}_α , with $\alpha = 1, 2$, have rather small couplings to the third pair of Higgs doublets H_u and H_d , i.e. $\lambda_{3\alpha} \approx 0$. This guarantees that \hat{S}_α develop rather small VEVs and the mixing between the neutral scalar components of \hat{S}_α , H_u and H_d can be negligibly small so that it can be even ignored in the leading approximation. In this context it is worth pointing out that if couplings κ_{i3} , $\lambda_{3\alpha}$, $\lambda_{\alpha 3}$ and λ_{33} are set to be small at the scale M_X then they will remain small at any scale below M_X .

Neglecting the Yukawa couplings $\lambda_{3\alpha}$ the low energy effective superpotential of the modified E₆SSM below the scale $\langle \hat{S}_3 \rangle$ can be written as

$$W_{eff} \simeq \lambda_{\alpha 1}\hat{S}_1(\hat{H}_\alpha^d\hat{H}_\alpha^u) + \kappa_{i1}\hat{S}_1(\hat{D}_i\hat{\bar{D}}_i) + \lambda_{\alpha 2}\hat{S}_2(\hat{H}_\alpha^d\hat{H}_\alpha^u) + \kappa_{i2}\hat{S}_2(\hat{D}_i\hat{\bar{D}}_i) \\ + \mu_{H_\alpha}(\hat{H}_\alpha^d\hat{H}_\alpha^u) + \mu_{D_i}(\hat{D}_i\hat{\bar{D}}_i) + W_{MSSM}(\mu \neq 0). \quad (10)$$

where $\alpha = 1, 2$ and $i = 1, 2, 3$. The superpotential (10) does not contain any mass terms that involve superfields \hat{S}_α . This implies that the fermion components of \hat{S}_α can be very light. In particular, the corresponding states can be lighter than 0.1 eV forming hot dark matter in the Universe. Such fermion states have negligible couplings to Z boson as well as other SM particles and therefore would not have been observed at earlier collider experiments. These states also do not change the branching ratios of the Z boson and Higgs decays[‡]. Moreover if Z' boson is sufficiently heavy the presence of such light fermion states does not affect Big Bang Nucleosynthesis [166].

The superpotential (10) contains ten new Yukawa couplings $\lambda_{\alpha 1}$, $\lambda_{\alpha 2}$, κ_{i1} and κ_{i2} . The running of these Yukawa couplings obey the following system of the renormalization group

[‡]The presence of very light neutral fermions in the particle spectrum might have interesting implications for the neutrino physics (see, for example [171]).

(RG) equations:

$$\begin{aligned}
\frac{d\lambda_{\alpha 1}}{dt} &= \frac{\lambda_{\alpha 1}}{(4\pi)^2} \left[2\lambda_{\alpha 1}^2 + 2\lambda_{\alpha 2}^2 + 2\left(\sum_{\beta} \lambda_{\beta 1}^2\right) + 3\left(\sum_j \kappa_{j1}^2\right) - 3g_2^2 \right. \\
&\quad \left. - \frac{3}{5}g_1^2 - \frac{19}{10}g_1'^2 \right] + \frac{\lambda_{\alpha 2}}{(4\pi)^2} \left[2\left(\sum_{\beta} \lambda_{\beta 1}\lambda_{\beta 2}\right) + 3\left(\sum_j \kappa_{j1}\kappa_{j2}\right) \right], \\
\frac{d\lambda_{\alpha 2}}{dt} &= \frac{\lambda_{\alpha 2}}{(4\pi)^2} \left[2\lambda_{\alpha 1}^2 + 2\lambda_{\alpha 2}^2 + 2\left(\sum_{\beta} \lambda_{\beta 2}^2\right) + 3\left(\sum_j \kappa_{j2}^2\right) - 3g_2^2 \right. \\
&\quad \left. - \frac{3}{5}g_1^2 - \frac{19}{10}g_1'^2 \right] + \frac{\lambda_{\alpha 1}}{(4\pi)^2} \left[2\left(\sum_{\beta} \lambda_{\beta 1}\lambda_{\beta 2}\right) + 3\left(\sum_j \kappa_{j1}\kappa_{j2}\right) \right], \\
\frac{d\kappa_{i1}}{dt} &= \frac{\kappa_{i1}}{(4\pi)^2} \left[2\kappa_{i1}^2 + 2\kappa_{i2}^2 + 2\left(\sum_{\beta} \lambda_{\beta 1}^2\right) + 3\left(\sum_j \kappa_{j1}^2\right) - \frac{16}{3}g_3^2 \right. \\
&\quad \left. - \frac{4}{15}g_1^2 - \frac{19}{10}g_1'^2 \right] + \frac{\kappa_{i2}}{(4\pi)^2} \left[2\left(\sum_{\beta} \lambda_{\beta 1}\lambda_{\beta 2}\right) + 3\left(\sum_j \kappa_{j1}\kappa_{j2}\right) \right], \\
\frac{d\kappa_{i2}}{dt} &= \frac{\kappa_{i2}}{(4\pi)^2} \left[2\kappa_{i1}^2 + 2\kappa_{i2}^2 + 2\left(\sum_{\beta} \lambda_{\beta 2}^2\right) + 3\left(\sum_j \kappa_{j2}^2\right) - \frac{16}{3}g_3^2 \right. \\
&\quad \left. - \frac{4}{15}g_1^2 - \frac{19}{10}g_1'^2 \right] + \frac{\kappa_{i1}}{(4\pi)^2} \left[2\left(\sum_{\beta} \lambda_{\beta 1}\lambda_{\beta 2}\right) + 3\left(\sum_j \kappa_{j1}\kappa_{j2}\right) \right].
\end{aligned} \tag{11}$$

The requirement of validity of perturbation theory up to the Grand Unification scale M_X restricts the interval of variations of these Yukawa couplings at low-energies. In our analysis here we use a set of one-loop RG equations (11) while the evolution of gauge couplings is calculated in the two-loop approximation.

3 750 GeV diphoton excess in the variant E_6 SSM

Turning now to a discussion of the 750 GeV diphoton excess recently observed by ATLAS and CMS in the framework of the variant of the E_6 SSM discussed in the previous section, whose effective superpotential is given by Eq. (10). This SUSY model involves two SM singlet superfields $\hat{S}_{1,2}$ plus a set of extra vector-like supermultiplets beyond the MSSM, including two pairs of inert Higgs doublets (\hat{H}_{α}^d and \hat{H}_{α}^u), as well as three generations of exotic quarks \hat{D}_i and $\overline{\hat{D}}_i$ with electric charges $\mp 1/3$.

The scenario discussed in this section is that the 750-760 GeV diphoton resonance may be identified as one or two scalars denoted $N_{1,2}$ and/or one or two pseudoscalars denoted $A_{1,2}$ contained in the two singlet superfields $\hat{S}_{1,2}$. The masses of these scalars and pseudoscalars arises from the soft SUSY breaking sector. However, to simplify our analysis, we assume that all other sparticles are sufficiently heavy so that their contributions to the production and decay rates of states with masses around 750 GeV can be ignored. Moreover the scenario under consideration implies that almost all exotic vector-like fermion mass states are heavier than 375 GeV so that the on-shell decays of N_{α} and A_{α} into the corresponding particles are not kinematically allowed.

Integrating out the heavy fermions corresponding to two pairs of inert Higgsino doublets \tilde{H}_{α}^d and \tilde{H}_{α}^u and three generations of vector-like D_i and \overline{D}_i fermions, which appear in the usual triangle loop diagrams, one obtains the effective Lagrangian which describes the interactions of

N_α and A_α with the SM gauge bosons,

$$\mathcal{L}_{eff} = \sum_\alpha \left(c_{1\alpha} N_\alpha B_{\mu\nu} B^{\mu\nu} + c_{2\alpha} N_\alpha W_{\mu\nu}^a W^{a\mu\nu} + c_{3\alpha} N_\alpha G_{\mu\nu}^\sigma G^{\sigma\mu\nu} \right. \\ \left. + \tilde{c}_{1\alpha} A_\alpha B_{\mu\nu} \tilde{B}^{\mu\nu} + \tilde{c}_{2\alpha} A_\alpha W_{\mu\nu}^a \tilde{W}^{a\mu\nu} + \tilde{c}_{3\alpha} A_\alpha G_{\mu\nu}^\sigma \tilde{G}^{\sigma\mu\nu} \right), \quad (12)$$

where

$$\begin{aligned} c_{1\alpha} &= \frac{\alpha_Y}{16\pi} \left[\sum_i \frac{2\kappa_{i\alpha}}{3\sqrt{2}\mu_{D_i}} A(x_{D_i}) + \sum_\beta \frac{\lambda_{\beta\alpha}}{\sqrt{2}\mu_{H_\beta}} A(x_{H_\beta}) \right], \\ c_{2\alpha} &= \frac{\alpha_2}{16\pi} \left[\sum_\beta \frac{\lambda_{\beta\alpha}}{\sqrt{2}\mu_{H_\beta}} A(x_{H_\beta}) \right], \\ c_{3\alpha} &= \frac{\alpha_3}{16\pi} \left[\sum_i \frac{\kappa_{i\alpha}}{\sqrt{2}\mu_{D_i}} A(x_{D_i}) \right], \end{aligned} \quad (13)$$

$$A(x) = 2x(1 + (1-x)\arcsin^2[1/\sqrt{x}]), \quad \text{for } x \geq 1.$$

In Eq. (12) $B_{\mu\nu}$, $W_{\mu\nu}^a$, $G_{\mu\nu}^\sigma$ are field strengths for the $U(1)_Y$, $SU(2)_W$ and $SU(3)_C$ gauge interactions respectively while $\tilde{G}^{\sigma\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\lambda\rho}G_{\lambda\rho}^\sigma$ etc. In Eqs. (13) $x_{D_i} = 4\mu_{D_i}^2/M_X^2$, $x_{H_\alpha} = 4\mu_{H_\alpha}^2/M_X^2$ and $\alpha_Y = 3\alpha_1/5$ whereas α_1 , α_2 and α_3 are (GUT normalised) gauge couplings of $U(1)_Y$, $SU(2)_W$ and $SU(3)_C$ interactions. In order to obtain analytic expressions for $\tilde{c}_{i\alpha}$ one should replace in Eqs. (13) $c_{i\alpha}$ by $\tilde{c}_{i\alpha}$ and substitute function $B(x)$ instead of $A(x)$, where

$$B(x) = 2x\arcsin^2[1/\sqrt{x}]. \quad (14)$$

Because in our analysis we focus on the diphoton decays of N_α and A_α that may lead to the 750 GeV diphoton excess it is convenient to use the effective Lagrangian that describes the interactions of these fields with the electromagnetic one. Using Eq. (12) one obtains

$$\mathcal{L}_{eff}^{\gamma\gamma} = \sum_\alpha \left(c_\alpha^\gamma N_\alpha F_{\mu\nu} F^{\mu\nu} + \tilde{c}_\alpha^\gamma A_\alpha F_{\mu\nu} \tilde{F}^{\mu\nu} \right), \quad (15)$$

where $c_\alpha^\gamma = c_{1\alpha}\cos^2\theta_W + c_{2\alpha}\sin^2\theta_W$, $\tilde{c}_\alpha^\gamma = \tilde{c}_{1\alpha}\cos^2\theta_W + \tilde{c}_{2\alpha}\sin^2\theta_W$ and $F_{\mu\nu}$ is a field strength associated with the electromagnetic interaction.

At the LHC the exotic states N_α and A_α can be predominantly produced through gluon fusion. When exotic quarks have masses below 1 TeV the corresponding production cross section is rather large and determined by the effective couplings $|c_{3\alpha}|^2$ and $|\tilde{c}_{3\alpha}|^2$. However such states mainly decay into a pair of gluons which is very problematic to detect at the LHC. Therefore possible collider signatures of these exotic states are associated with their decays into WW , ZZ , γZ and $\gamma\gamma$. Since W and Z decay mostly into quarks the process $pp \rightarrow N_\alpha(A_\alpha) \rightarrow \gamma\gamma$ tends to be one of the most promising channels to search for such resonances. In the limit when exotic states decay predominantly into a pair of gluons the branching ratios of $N_\alpha \rightarrow \gamma\gamma$ and $A_\alpha \rightarrow \gamma\gamma$ are proportional to $|c_\alpha^\gamma|^2/|c_{3\alpha}|^2$ and $|\tilde{c}_\alpha^\gamma|^2/|\tilde{c}_{3\alpha}|^2$ respectively. As a consequence cross sections $\sigma(pp \rightarrow N_\alpha(A_\alpha) \rightarrow \gamma\gamma)$ do not depend on $|c_{3\alpha}|^2$ and $|\tilde{c}_{3\alpha}|^2$. The corresponding signal strengths are basically defined by the partial decay widths $\Gamma(N_\alpha \rightarrow \gamma\gamma)$ and $\Gamma(A_\alpha \rightarrow \gamma\gamma)$.

The cross sections of the processes that may result in the 750 GeV diphoton excess can be written as

$$\sigma(pp \rightarrow X_\alpha \rightarrow \gamma\gamma) \simeq \frac{C_{gg}}{M_{X_\alpha} s \Gamma_{X_\alpha}} \Gamma(X_\alpha \rightarrow gg) \Gamma(X_\alpha \rightarrow \gamma\gamma), \quad (16)$$

where X_α is either N_α or A_α exotic states, Γ_{X_α} is a total decay width of the resonance X_α while $C_{gg} \simeq 3163$, $\sqrt{s} \simeq 13$ TeV and M_{X_α} is the mass of the appropriate exotic state which should be somewhat around 750 GeV. The partial decay widths of the corresponding resonances are given by

$$\begin{aligned}\Gamma(N_\alpha \rightarrow gg) &= \frac{2}{\pi} M_{N_\alpha}^3 |c_{3\alpha}|^2, & \Gamma(A_\alpha \rightarrow gg) &= \frac{2}{\pi} M_{A_\alpha}^3 |\tilde{c}_{3\alpha}|^2, \\ \Gamma(N_\alpha \rightarrow \gamma\gamma) &= \frac{M_{N_\alpha}^3 |c_\alpha^\gamma|^2}{4\pi}, & \Gamma(A_\alpha \rightarrow \gamma\gamma) &= \frac{M_{A_\alpha}^3 |\tilde{c}_\alpha^\gamma|^2}{4\pi}.\end{aligned}\tag{17}$$

In the limit when $\Gamma_{X_\alpha} \approx \Gamma(X_\alpha \rightarrow gg)$ the dependence of the cross section (16) on $\Gamma(X_\alpha \rightarrow gg)$ disappear and its value is determined by the partial decay width $\Gamma(X_\alpha \rightarrow \gamma\gamma)$ as one could naively expect. In this case, as it was pointed out in [10], one can obtain $\sigma(pp \rightarrow \gamma\gamma) \approx 8$ fb at the 13 TeV LHC if

$$\frac{\Gamma(X_\alpha \rightarrow \gamma\gamma)}{M_{X_\alpha}} = 1.1 \times 10^{-6}.\tag{18}$$

Then the cross section $\sigma_{\gamma\gamma} \approx \sigma(pp \rightarrow \gamma\gamma)$ for arbitrary partial decay widths of $X_\alpha \rightarrow \gamma\gamma$ can be approximately estimated as

$$\sigma_{\gamma\gamma} \simeq 7.3 \text{ fb} \times \text{BR}(X_\alpha \rightarrow gg) \times \left(\frac{\Gamma(X_\alpha \rightarrow \gamma\gamma)}{M_{X_\alpha}} \times 10^6 \right).\tag{19}$$

where the branching ratios associated with the decays of exotic states into gluons g and vector bosons V ($V = \gamma, W^\pm, Z$) are given by

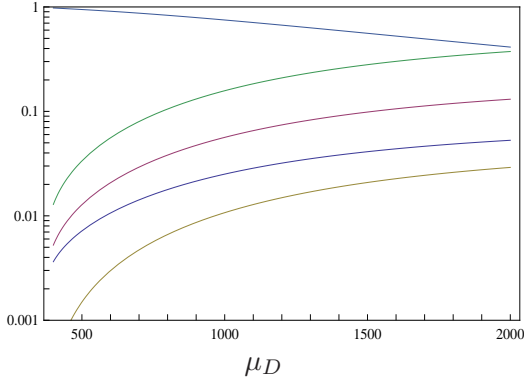
$$\text{BR}(X_\alpha \rightarrow gg) = \frac{\Gamma(X_\alpha \rightarrow gg)}{\Gamma_{X_\alpha}}, \quad \text{BR}(X_\alpha \rightarrow VV) = \frac{\Gamma(X_\alpha \rightarrow VV)}{\Gamma_{X_\alpha}}.\tag{20}$$

In Eqs. (20) $\Gamma(X_\alpha \rightarrow gg)$ and $\Gamma(X_\alpha \rightarrow VV)$ are partial decay widths that correspond to the exotic state decays into a pair of gluons and a pair of vector bosons respectively whereas Γ_{X_α} is a total decay width of this state.

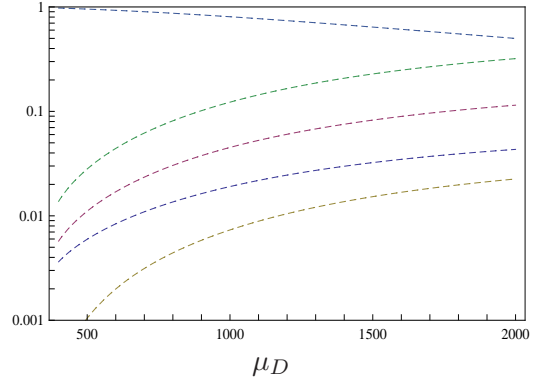
3.1 One scalar/pseudoscalar case

Let us now consider the scenario when one of the scalar/pseudoscalar exotic states (N_1 or A_1) has a mass which is rather close to 750 GeV. From Eqs. (13)–(15) and (17) it follows that the diphoton decay rates of these new bosons and the corresponding signal strength depend very strongly on the values of the Yukawa couplings $\lambda_{\alpha 1}$ and κ_{i1} . On the other hand the growth of these Yukawa couplings at low energies entails the increase of their values at the Grand Unification scale M_X resulting in the appearance of the Landau pole that spoils the applicability of perturbation theory at high energies (see, for example [172]). The requirement of validity of perturbation theory up to the scale M_X sets an upper bound on the low energy value of $\lambda_{\alpha 1}$ and κ_i . In our analysis we use two-loop SM RG equations to compute the values of the gauge couplings at the scale $Q = 2$ TeV. Above this scale we use two-loop RG equations for the gauge couplings and one-loop RG equations for the Yukawa couplings including the ones given by Eq. (11) to analyse the RG flow of these couplings. In the simplest case when $\lambda_{\alpha 1} = \kappa_{i1}$ our numerical analysis indicates that the values of these couplings at the scale $Q = 2$ TeV should not exceed 0.6.

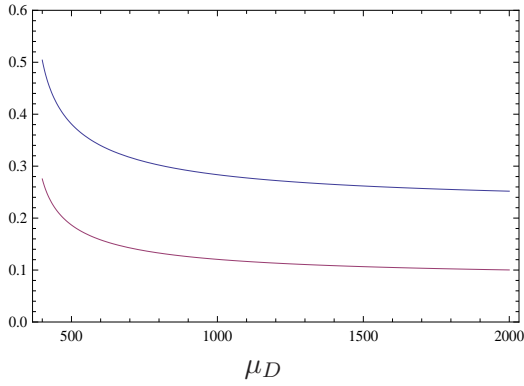
The upper bound on the coupling $\lambda_{\alpha 1}$ becomes less stringent when κ_{i1} are small. In the limit when all κ_{i1} vanish the value of $\lambda_{11} = \lambda_{21}$ has to remain smaller than 0.81 to ensure the applicability of perturbation theory up to the GUT scale. Although in this case $\Gamma(A_1 \rightarrow \gamma\gamma)$ and $\Gamma(N_1 \rightarrow \gamma\gamma)$ attain their maximal value the production cross sections of exotic states N_1 or A_1 are negligibly small since they are determined by the low-energy values of κ_{i1} . The upper

$\text{BR}(A_1 \rightarrow gg, WW, ZZ, \gamma\gamma, \gamma Z)$


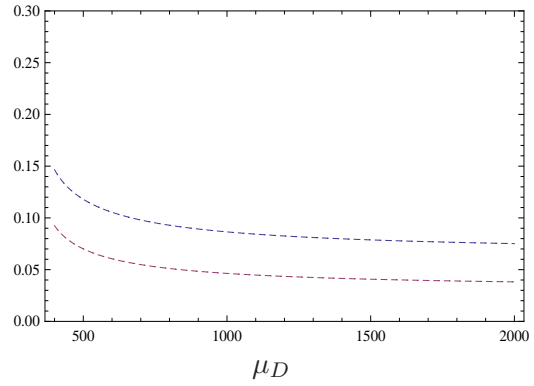
(a)

 $\text{BR}(N_1 \rightarrow gg, WW, ZZ, \gamma\gamma, \gamma Z)$


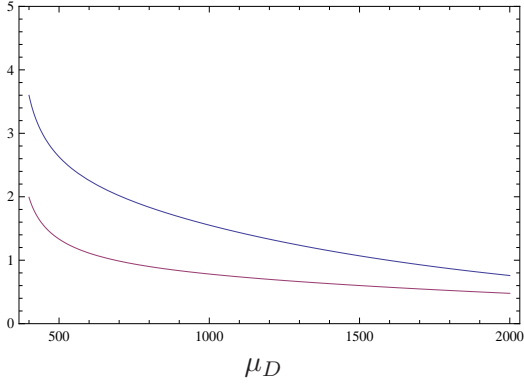
(b)

 $\frac{\Gamma(A_1 \rightarrow \gamma\gamma)}{M_X} \times 10^6$


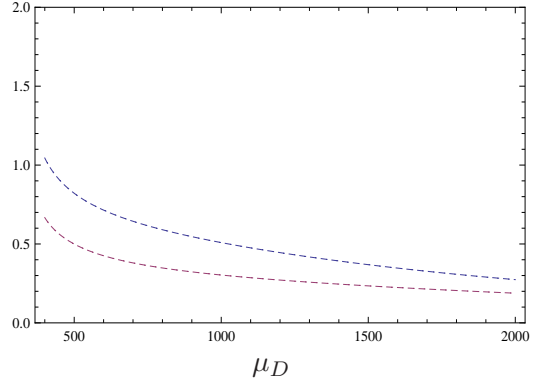
(c)

 $\frac{\Gamma(N_1 \rightarrow \gamma\gamma)}{M_X} \times 10^6$


(d)

 $\sigma(pp \rightarrow A_1 \rightarrow \gamma\gamma)[\text{fb}]$


(e)

 $\sigma(pp \rightarrow N_1 \rightarrow \gamma\gamma)[\text{fb}]$


(f)

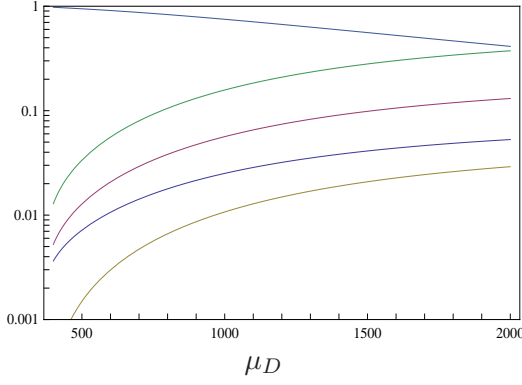
Figure 1: Predictions for the one pseudoscalar (left panels) or one scalar (right panels) case. In all cases the masses of vector-like quarks are set to be equal, i.e. $\mu_{D_i} = \mu_D$, whereas $\lambda_{\alpha 1} = \kappa_{i1} = 0.6$, $\lambda_{\alpha 2} = \kappa_{i2} = 0$. In (a) the branching ratios of the decays of A_1 into γZ (lowest solid line), $\gamma\gamma$ (second lowest solid line), ZZ (third lowest solid line), WW (second highest solid line) and gg (highest solid line) as a function of exotic quark masses μ_D for $M_{A_1} \simeq 750$ GeV. In (b) the branching ratios of the decays of N_1 into γZ (lowest dashed line), $\gamma\gamma$ (second lowest dashed line), ZZ (third lowest dashed line), WW (second highest dashed line) and gg (highest dashed line) as a function of μ_D for $M_{N_1} \simeq 750$ GeV. In (c) the ratios $\Gamma(A_1 \rightarrow \gamma\gamma)/M_X$ as a function of μ_D for $M_{A_1} \simeq 750$ GeV. The upper and lower solid lines correspond to the scenarios with $\mu_{H_\alpha} = 400$ GeV and $\mu_{H_\alpha} = 500$ GeV. In (d) the ratios $\Gamma(N_1 \rightarrow \gamma\gamma)/M_X$ as a function of μ_D for $M_{N_1} \simeq 750$ GeV. The upper and lower dashed lines correspond to the scenarios with $\mu_{H_\alpha} = 400$ GeV and $\mu_{H_\alpha} = 500$ GeV. In (e) the cross sections $\sigma(pp \rightarrow A_1 \rightarrow \gamma\gamma)$ in fb as a function of μ_D for $M_{A_1} \simeq 750$ GeV. The upper and lower solid lines represent the scenarios with $\mu_{H_\alpha} = 400$ GeV and $\mu_{H_\alpha} = 500$ GeV. In (f) the cross sections $\sigma(pp \rightarrow N_1 \rightarrow \gamma\gamma)$ in fb as a function of μ_D for $M_{N_1} \simeq 750$ GeV. The upper and lower dashed lines represent the scenarios with $\mu_{H_\alpha} = 400$ GeV and $\mu_{H_\alpha} = 500$ GeV.

bounds on κ_{i1} can be also significantly relaxed when $\lambda_{11} = \lambda_{21} = 0$. If this is a case then the requirement of the validity of perturbation theory implies that $\kappa_{11} = \kappa_{21} = \kappa_{31} \lesssim 0.79$. However in this limit the diphoton production rate associated with the presence A_1 or N_1 is again negligibly small because the corresponding partial decay width vanish. Thus in this section we focus on the scenario with $\lambda_{\alpha 1} = \kappa_{i1} = 0.6$. This choice of parameters guarantees that the production cross sections of N_1 and A_1 as well as their partial decay width can be sufficiently large.

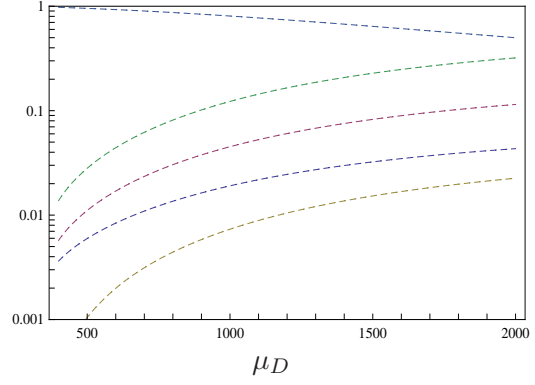
In Figs. 1a and 1b the dependence of the branching ratios of the exotic pseudoscalar and scalar states on the masses of exotic quarks is examined. To simplify our analysis the masses of all exotic quarks are set to be equal while the masses of all inert Higgsinos are assumed to be around 400 GeV. From Fig. 1a and 1b it follows that the exotic pseudoscalar and scalar states decay predominantly into a pair of gluons when the masses of exotic quarks μ_D are below 1 TeV. Moreover if μ_D is close to 400 – 500 GeV all other branching ratios are negligibly small. With increasing μ_D the branching ratio of the exotic pseudoscalar (scalar) state decays into gluons decreases whereas the branching ratios of the decays of this state into W^+W^- , ZZ , $\gamma\gamma$ and γZ increase. The branching ratios of $A_1(N_1) \rightarrow WW$ and $A_1(N_1) \rightarrow ZZ$ are the second and third largest ones. The branching ratio of $A_1(N_1) \rightarrow \gamma\gamma$ is considerably smaller but still larger than $A_1(N_1) \rightarrow \gamma Z$. Although the branching ratios of $A_1(N_1) \rightarrow WW$ and $A_1(N_1) \rightarrow ZZ$ can be a substantially bigger than the branching ratio $A_1(N_1) \rightarrow \gamma\gamma$ their experimental detection is more problematic because W and Z decays mainly into quarks. When μ_D is around 1 TeV the branching ratio of $A_1(N_1) \rightarrow gg$ is still the largest one and constitutes about 75%(80%) while for $\mu_D \simeq 2$ TeV the branching ratios of $A_1(N_1) \rightarrow gg$ and $A_1(N_1) \rightarrow WW$ become comparable.

In Fig. 1c and 1d we explore the dependence of the partial decay widths associated with the decays of the exotic pseudoscalar and scalar states into a pair of photons on the masses of exotic quarks and inert Higgsinos μ_{H_α} . One can see that these decay widths decrease very rapidly with increasing μ_{H_α} . The dependence on the masses of exotic quarks is weaker because these states carry small electric charges $\pm 1/3$. Since here we assume that κ_{i1}/μ_{D_i} and $\lambda_{\alpha 1}/\mu_{H_\alpha}$ have the same sign the growth of either exotic quark masses or μ_{H_α} results in the reduction of the corresponding decay rate. When μ_D is larger than 1.5 TeV the dependence of the partial decay widths under consideration becomes rather weak. From Fig. 1c and 1d it is easy to see that the partial width of the decays $A_1 \rightarrow \gamma\gamma$ is substantially larger than the one associated with $N_1 \rightarrow \gamma\gamma$ leading to the larger value of the cross sections $\sigma(pp \rightarrow A_1 \rightarrow \gamma\gamma)$ as compared with $\sigma(pp \rightarrow N_1 \rightarrow \gamma\gamma)$.

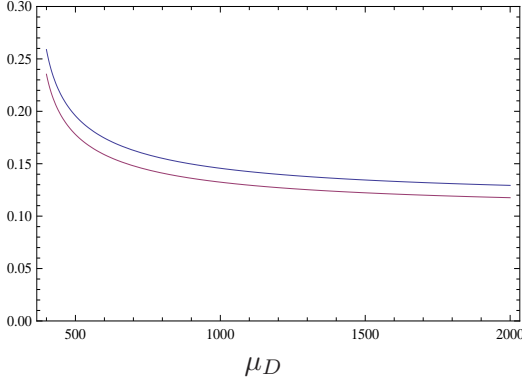
In our analysis we use Eq. (19) to estimate the values of the cross sections $\sigma(pp \rightarrow A_1 \rightarrow \gamma\gamma)$ and $\sigma(pp \rightarrow N_1 \rightarrow \gamma\gamma)$ at the 13 TeV LHC. The results of our investigation are shown in Figs. 1e and 1f. In the case of scalar exotic states with mass 750 GeV this cross section tends to be substantially smaller than 1 fb. The presence of 750 GeV exotic pseudoscalar can lead to the considerably stronger signal in the diphoton channel. When all exotic quarks have masses around 400–500 GeV the corresponding cross section can reach 2–3 fb. Somewhat stronger signal can be obtained if we assume that both scalar and pseudoscalar exotic states have masses which are close to 750 GeV. In this case the sum of the cross sections $\sigma(pp \rightarrow A_1 \rightarrow \gamma\gamma) + \sigma(pp \rightarrow N_1 \rightarrow \gamma\gamma)$ can reach 4.5 fb if exotic quarks have masses about 400 GeV. The existence of two nearly degenerate resonances may also explain why the analysis performed by the ATLAS collaboration leads to the relatively large best-fit width which is about 45 GeV. Unfortunately, the cross sections mentioned above decreases substantially with increasing exotic quark masses. Indeed, if $\mu_D \gtrsim 1$ TeV the sum of the cross sections $\sigma(pp \rightarrow A_1 \rightarrow \gamma\gamma) + \sigma(pp \rightarrow N_1 \rightarrow \gamma\gamma)$ does not exceed 2 fb. These cross sections continue to fall even for $\mu_D \gtrsim 1.5$ TeV when the corresponding partial decay widths are rather close to their lower saturation limits because the branching ratios associated with the decays of A_1 and N_1 into a pair of gluons decrease with increasing μ_D .

$\text{BR}(A_{1,2} \rightarrow gg, WW, ZZ, \gamma\gamma, \gamma Z)$


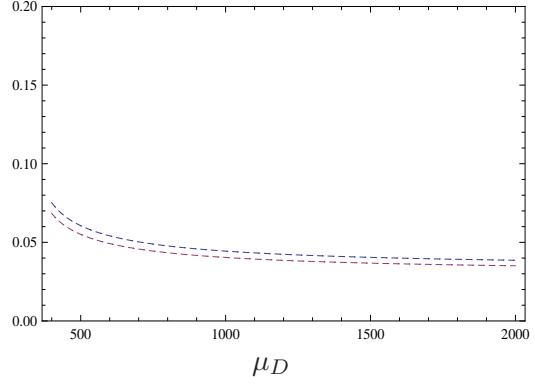
(a)

 $\text{BR}(N_{1,2} \rightarrow gg, WW, ZZ, \gamma\gamma, \gamma Z)$


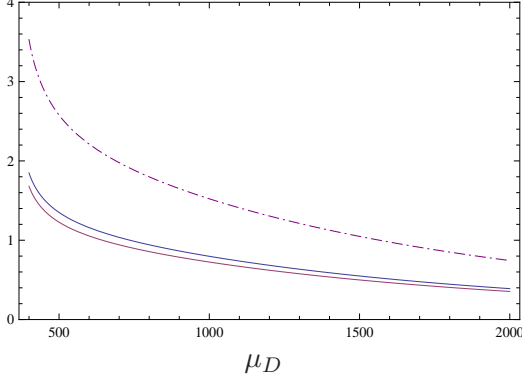
(b)

 $\frac{\Gamma(A_{1,2} \rightarrow \gamma\gamma)}{M_X} \times 10^6$


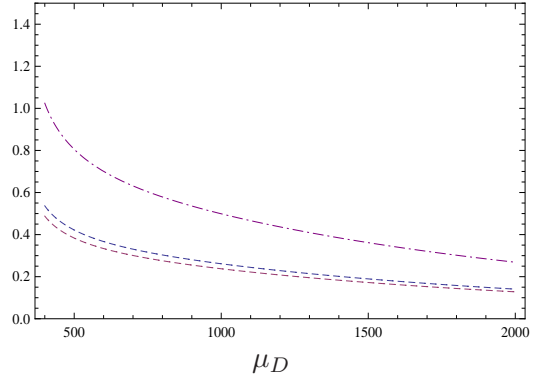
(c)

 $\frac{\Gamma(N_{1,2} \rightarrow \gamma\gamma)}{M_X} \times 10^6$


(d)

 $\sigma(pp \rightarrow A_{1,2} \rightarrow \gamma\gamma)[\text{fb}]$


(e)

 $\sigma(pp \rightarrow N_{1,2} \rightarrow \gamma\gamma)[\text{fb}]$


(f)

Figure 2: Predictions for two degenerate pseudoscalars (left panels) or two degenerate scalars (right panels) case. In all cases $\mu_{D_i} = \mu_D$, while $\mu_{H_\alpha} = 400$ GeV, $\lambda_{\alpha 1} = \kappa_{i1} = 0.43$ and $\lambda_{\alpha 2} = \kappa_{i2} = 0.41$. In (a) the branching ratios of the decays of $A_{1,2}$ into γZ (lowest solid line), $\gamma\gamma$ (second lowest solid line), ZZ (third lowest solid line), WW (second highest solid line) and gg (highest solid line) as a function of exotic quark masses μ_D for $M_{A_{1,2}} \simeq 750$ GeV. In (b) the branching ratios of the decays of $N_{1,2}$ into γZ (lowest dashed line), $\gamma\gamma$ (second lowest dashed line), ZZ (third lowest dashed line), WW (second highest dashed line) and gg (highest dashed line) as a function of μ_D for $M_{N_{1,2}} \simeq 750$ GeV. In (c) the ratios $\Gamma(A_1 \rightarrow \gamma\gamma)/M_X$ (upper solid line) and $\Gamma(A_2 \rightarrow \gamma\gamma)/M_X$ (lower solid line) as a function of μ_D for $M_{A_{1,2}} \simeq 750$ GeV. In (d) the ratios $\Gamma(N_1 \rightarrow \gamma\gamma)/M_X$ (upper dashed line) and $\Gamma(N_2 \rightarrow \gamma\gamma)/M_X$ (lower dashed line) as a function of μ_D for $M_{N_{1,2}} \simeq 750$ GeV. In (e) the cross sections (fb) $\sigma(pp \rightarrow A_1 \rightarrow \gamma\gamma)$ (upper solid line) and $\sigma(pp \rightarrow A_2 \rightarrow \gamma\gamma)$ (lower solid line) as a function of μ_D for $M_{A_{1,2}} \simeq 750$ GeV. The dashed-dotted line correspond to the sum of these cross sections. In (f) the cross sections (fb) $\sigma(pp \rightarrow N_1 \rightarrow \gamma\gamma)$ (upper dashed line) and $\sigma(pp \rightarrow N_2 \rightarrow \gamma\gamma)$ (lower dashed line) as a function of μ_D for $M_{N_{1,2}} \simeq 750$ GeV. The dashed-dotted line correspond to the sum of these cross sections.

3.2 Two degenerate scalar/pseudoscalar case

Now let us assume that there are two superfields \hat{S}_1 and \hat{S}_2 that have sufficiently large Yukawa couplings to the exotic quark and inert Higgsino states and can contribute to the measured cross section $pp \rightarrow \gamma\gamma$. In other words we assume that scalar and pseudoscalar components of both superfields can have masses around 750 GeV. Naively one may expect that this could allow to enhance the theoretical prediction for the cross section $pp \rightarrow \gamma\gamma$. Again we start from the simplest case when all Yukawa couplings are the same. Then the numerical analysis indicates that in this case the requirement of the validity of perturbation theory up to the scale M_X sets even more stringent upper bound on the low energy value of the Yukawa couplings as compared with the one scalar/pseudoscalar case. Indeed, using the one-loop RG equations (11) and two-loop RG equations for the gauge couplings one obtains that $\lambda_{\alpha 1} = \kappa_{i1} = \lambda_{\alpha 2} = \kappa_{i2} = \lambda_0 \lesssim 0.43$. Smaller values of the Yukawa couplings do not affect the branching ratios of A_1 and N_1 . Moreover A_2 and A_1 as well as N_2 and N_1 have basically the same branching ratios. This is because partial decay widths of $A_{1,2}$ and $N_{1,2}$ as well as the corresponding total widths are proportional to λ_0^2 . As a consequence in the leading approximation branching ratios do not depend on λ_0 (see Fig. 2a and 2b). On the other hand as one can see from Fig. 1c, 1d, 1e and 1f the partial decay widths of $A_{1,2} \rightarrow \gamma\gamma$ and $N_{1,2} \rightarrow \gamma\gamma$ as well as the cross sections $\sigma(pp \rightarrow A_{1,2} \rightarrow \gamma\gamma)$ and $\sigma(pp \rightarrow N_{1,2} \rightarrow \gamma\gamma)$ are reduced by factor 2 because of the smaller values of the Yukawa couplings. If all exotic states A_1 and A_2 as well as N_1 and N_2 are nearly degenerate around 750 GeV so that their distinction is not possible within present experimental accuracy, then the superpositions of rates from these bosons basically reproduces the corresponding rates in the one scalar/pseudoscalar case (see Figs. 1e, 1f, 2e and 2f). Thus, it seems rather problematic to achieve any enhancement of the signal in the diphoton channel in the scenario when all Yukawa couplings are equal or reasonably close to each other.

3.3 Maximal mixing scenario

Following on from the discussion in the previous subsection, there is one case when a modest enhancement of the signal in the diphoton channel can be achieved. This happens in the so-called maximal mixing scenario when the masses of exotic scalars as well as the masses of exotic pseudoscalars are rather close to 750 GeV and the breakdown of SUSY gives rise to the mixing of these states preserving CP conservation. In this case one can expect that the mixing angles between CP-odd exotic states and CP-even exotic states tend to be rather large, i.e. about $\pm\pi/4$, because these bosons are nearly degenerate. To simplify our analysis here we set these angles to be equal to $\pi/4$. Then the scalar components of the superfields S_1 and S_2 can be expressed in terms of the mass eigenstates N_1 , N_2 , A_1 and A_2 as follows

$$S_1 = \frac{1}{2} (N_1 + N_2 + i(A_1 + A_2)) , \quad S_2 = \frac{1}{2} (N_1 - N_2 + i(A_1 - A_2)) . \quad (21)$$

In addition we assume that only superfield S_1 couples to the inert Higgsino states, i.e. $\lambda_{\alpha 2} = 0$, and only superfield S_2 couples to the exotic quarks, i.e. $\kappa_{i1} = 0$. In this limit the requirement of the validity of perturbation theory up to the scale M_X implies that $\lambda_{\alpha 1} = \lambda_0 \lesssim 0.8$ and $\kappa_{i2} = \kappa_0 \lesssim 0.79$.

Setting $\mu_{H_\alpha} = \mu_H$, $\mu_{D_i} = \mu_D$ and $M_{N_1} \simeq M_{N_2} \simeq M_{A_1} \simeq M_{A_2} \simeq M_X = 750$ GeV one can obtain simple analytical expressions for the partial decay widths of N_1 , A_1 , N_2 and A_2 into a

pair of photons

$$\Gamma(N_1 \rightarrow \gamma\gamma) = \frac{\alpha^2 M_X^3}{256\pi^3} \left| \frac{\lambda_0}{\mu_H} A(x_H) + \frac{\kappa_0}{2\mu_D} A(x_D) \right|^2, \quad (22)$$

$$\Gamma(A_1 \rightarrow \gamma\gamma) = \frac{\alpha^2 M_X^3}{256\pi^3} \left| \frac{\lambda_0}{\mu_H} B(x_H) + \frac{\kappa_0}{2\mu_D} B(x_D) \right|^2, \quad (23)$$

$$\Gamma(N_2 \rightarrow \gamma\gamma) = \frac{\alpha^2 M_X^3}{256\pi^3} \left| \frac{\lambda_0}{\mu_H} A(x_H) - \frac{\kappa_0}{2\mu_D} A(x_D) \right|^2, \quad (24)$$

$$\Gamma(A_2 \rightarrow \gamma\gamma) = \frac{\alpha^2 M_X^3}{256\pi^3} \left| \frac{\lambda_0}{\mu_H} B(x_H) - \frac{\kappa_0}{2\mu_D} B(x_D) \right|^2, \quad (25)$$

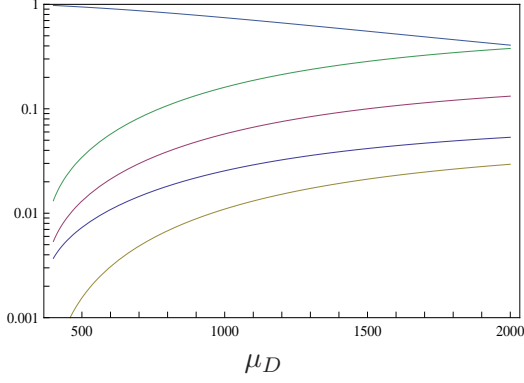
where $x_D = 4\mu_D^2/M_X^2$ and $x_H = 4\mu_H^2/M_X^2$. Assuming, that κ_0/μ_D and λ_0/μ_D have the same sign, Eqs. (22) and (23) are very similar to the ones which was used before for the calculation of the corresponding partial decay widths in one scalar/pseudoscalar case. Because the expressions for other partial decay widths are also very similar the branching ratios shown in Figs. 3a and 3b are almost the same as in Figs. 1a and 1b. At the same time in the case of N_2 and A_2 destructive interference between the contributions of exotic quarks and inert Higgsinos occurs. This leads to the suppression of the diphoton partial decay width. As a consequence when exotic quarks are lighter than 1 TeV the branching ratios of the decays $N_2 \rightarrow \gamma\gamma$ and $A_2 \rightarrow \gamma\gamma$ are the lowest ones (see Figs. 3c and 3d).

As before from Fig. 3 it follows that all exotic states N_1 , A_1 , N_2 and A_2 decay mainly into a pair of gluons. The corresponding branching ratio decreases with increasing μ_D because $c_{3\alpha}$ and $\tilde{c}_{3\alpha}$ diminish. The branching ratios of the decay of these states into WW and ZZ are the second largest and third largest ones. These branching ratios are substantially larger than the ones associated with the decays of exotic states into $\gamma\gamma$ and γZ . In the case of N_2 and A_2 the branching ratios of the decay of these states into WW can be an order of magnitude larger than the branching ratios of $N_2 \rightarrow \gamma\gamma$ and $A_2 \rightarrow \gamma\gamma$. Nevertheless the observation of the decays of N_α and A_α into pairs of WW and ZZ tend to be more problematic since W and Z decay mostly into quarks. All branching ratios of the exotic scalar and pseudoscalar decays except the largest one grow with increasing μ_D . As a result for $\mu_D \simeq 2$ TeV the branching ratios of $A_\alpha(N_\alpha) \rightarrow gg$ and $A_\alpha(N_\alpha) \rightarrow WW$ become sufficiently close.

The dependence of the partial decay widths and the corresponding cross sections at the 13 TeV LHC associated with the decays of the exotic pseudoscalar and scalar states into a pair of photons on the exotic quark masses is shown in Fig. 4. The results of our calculations for N_1 and A_1 are very similar to the ones obtained in the one scalar/pseudoscalar case (see Fig. 2e and 2f). The partial decay widths and the cross sections $\sigma(pp \rightarrow A_1(N_1) \rightarrow \gamma\gamma)$ are just a bit smaller since the Yukawa couplings of A_1 and N_1 to the exotic quarks and inert Higgsino states are slightly smaller. They decrease with increasing the masses of exotic quarks μ_D as before. On the contrary, the partial decay widths of $N_2 \rightarrow \gamma\gamma$ and $A_2 \rightarrow \gamma\gamma$ increase with increasing the exotic quark masses for fixed values of inert Higgsino masses because of the destructive interference mentioned above. They attain their maximal values for $\mu_D \gg 1$ TeV when the contribution of the exotic quarks to the partial decay widths become vanishingly small. The cross sections $\sigma(pp \rightarrow A_2(N_2) \rightarrow \gamma\gamma)$ also increase with increasing exotic quark masses when $\mu_D \lesssim 700$ GeV. However if exotic quarks are considerably heavier than 1 TeV then these cross sections become smaller for larger μ_D since the branching ratios of $A_2(N_2) \rightarrow gg$ diminish.

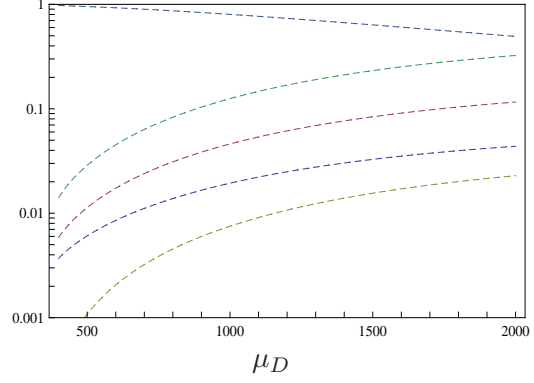
The sums of the cross sections $\sigma(pp \rightarrow N_1 \rightarrow \gamma\gamma) + \sigma(pp \rightarrow N_2 \rightarrow \gamma\gamma)$ and $\sigma(pp \rightarrow A_1 \rightarrow \gamma\gamma) + \sigma(pp \rightarrow A_2 \rightarrow \gamma\gamma)$ that correspond to the case when all exotic scalar and pseudoscalar states have

$\text{BR}(A_1 \rightarrow gg, WW, ZZ, \gamma\gamma, \gamma Z)$



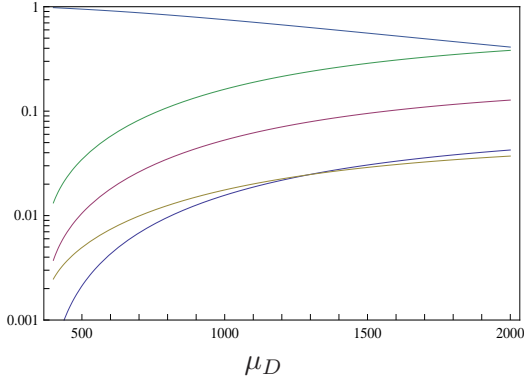
(a)

$\text{BR}(N_1 \rightarrow gg, WW, ZZ, \gamma\gamma, \gamma Z)$



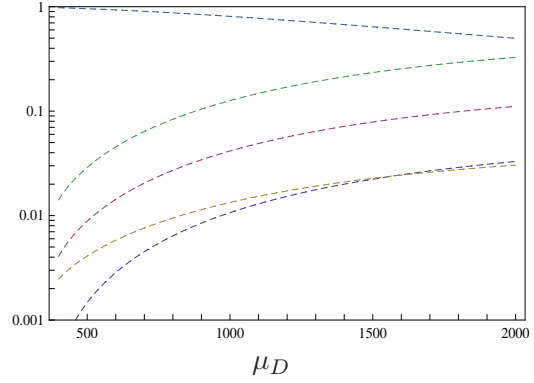
(b)

$\text{BR}(A_2 \rightarrow gg, WW, ZZ, \gamma Z, \gamma\gamma)$



(c)

$\text{BR}(N_2 \rightarrow gg, WW, ZZ, \gamma Z, \gamma\gamma)$



(d)

Figure 3: Predictions for the branching ratios in the maximal mixing scenario for two maximally mixed pseudoscalars (left panels) or two maximally mixed scalars (right panels) case. In all cases the masses of exotic quarks are set to be equal, i.e. $\mu_{D_i} = \mu_D$, while $\mu_{H_\alpha} = 400$ GeV, $\lambda_{\alpha 1} = 0.8$, $\kappa_{i2} = 0.79$ and $\kappa_{i1} = \lambda_{\alpha 2} = 0$. In (a) the branching ratios of the decays of A_1 into γZ (lowest solid line), $\gamma\gamma$ (second lowest solid line), ZZ (third lowest solid line), WW (second highest solid line) and gg (highest solid line) as a function of exotic quark masses for $M_{A_1} \simeq 750$ GeV. In (b) the branching ratios of the decays of N_1 into γZ (lowest dashed line), $\gamma\gamma$ (second lowest dashed line), ZZ (third lowest dashed line), WW (second highest dashed line) and gg (highest dashed line) as a function of exotic quark masses for $M_{N_1} \simeq 750$ GeV. In (c) the branching ratios of the decays of A_2 into $\gamma\gamma$ (lowest solid line), γZ (second lowest solid line), ZZ (third lowest solid line), WW (second highest solid line) and gg (highest solid line) as a function of exotic quark masses for $M_{A_1} \simeq 750$ GeV. In (d) the branching ratios of the decays of N_2 into $\gamma\gamma$ (lowest dashed line), γZ (second lowest dashed line), ZZ (third lowest dashed line), WW (second highest dashed line) and gg (highest dashed line) as a function of exotic quark masses for $M_{N_1} \simeq 750$ GeV.

masses around 750 GeV decreases with increasing μ_D (see Figs. 4c and 4d). At large values of the exotic quark masses these cross sections are bigger than the ones in the one scalar/pseudoscalar case shown in Fig. 1e and 1f. This is because the requirement of the validity of perturbation theory up to the scale M_X allows for larger values of $\lambda_{\alpha 1}$ in the maximal mixing scenario as compared with the one scalar/pseudoscalar case. From Figs. 4c and 4d one can see that the sum of all cross section that includes contributions of all scalar and pseudoscalar states with masses around 750 GeV changes from 4.5 fb to 3 fb when the exotic quark masses vary from 400 GeV to 1 TeV. The presence of such nearly degenerate states in the particle spectrum may also provide an explanation why the value of the best-fit width of the resonance obtained by ATLAS collaboration is so large.

4 Conclusions

In this paper we have proposed a variant of the E_6 SSM in which the third singlet S_3 breaks the gauged $U(1)_N$ above the TeV scale, which predicts a Z'_N , vector-like colour triplet and charge $\mp 1/3$ quarks D, \bar{D} , and two families of inert Higgsinos, all of which should be observed at LHC Run 2, plus the two lighter singlets $\hat{S}_{1,2}$ with masses around 750 GeV which are candidates for the recently observed diphoton excess. We have calculated the branching ratios and cross-sections for the two scalars $N_{1,2}$ and two pseudoscalars $A_{1,2}$ associated with $\hat{S}_{1,2}$, including possible degeneracies and maximal mixing, subject to the constraint that their couplings remain perturbative up to the unification scale.

Our results show that this variant of the E_6 SSM with two nearly degenerate pseudoscalars $A_{1,2}$ with masses around 750 GeV, may give rise to cross sections of $pp \rightarrow \gamma\gamma$ that can be as large as about 3 fb providing that the inert Higgsino states have masses around 400 GeV, while the three generations of D, \bar{D} are lighter than about 1 TeV. If the two nearly degenerate scalars $N_{1,2}$ also have masses around 750 GeV, then these cross-sections may be further boosted by about 1 fb, assuming that they are at present unresolvable. The existence of nearly degenerate spinless singlets provides an explanation for why the best-fit width of the 750 GeV resonance obtained by the ATLAS collaboration is apparently so large, i.e. about 45 GeV. However further data from Run 2 should begin to resolve the two separate pseudoscalar states $A_{1,2}$ (plus perhaps the two scalar states $N_{1,2}$).

Finally we emphasise that the three families of light vector-like D-quarks around 1 TeV and two families of inert Higgsinos around 400 GeV, although not currently ruled out because of their non-standard decay patterns, should be observable in dedicated searches at Run 2 of the LHC. The Z'_N gauge boson also remains a prediction of the E_6 SSM. In addition, the proposed variant E_6 SSM also predicts further decay modes of the 750 GeV resonance into WW , ZZ and γZ that might be possible to observe in the Run 2 at the LHC.

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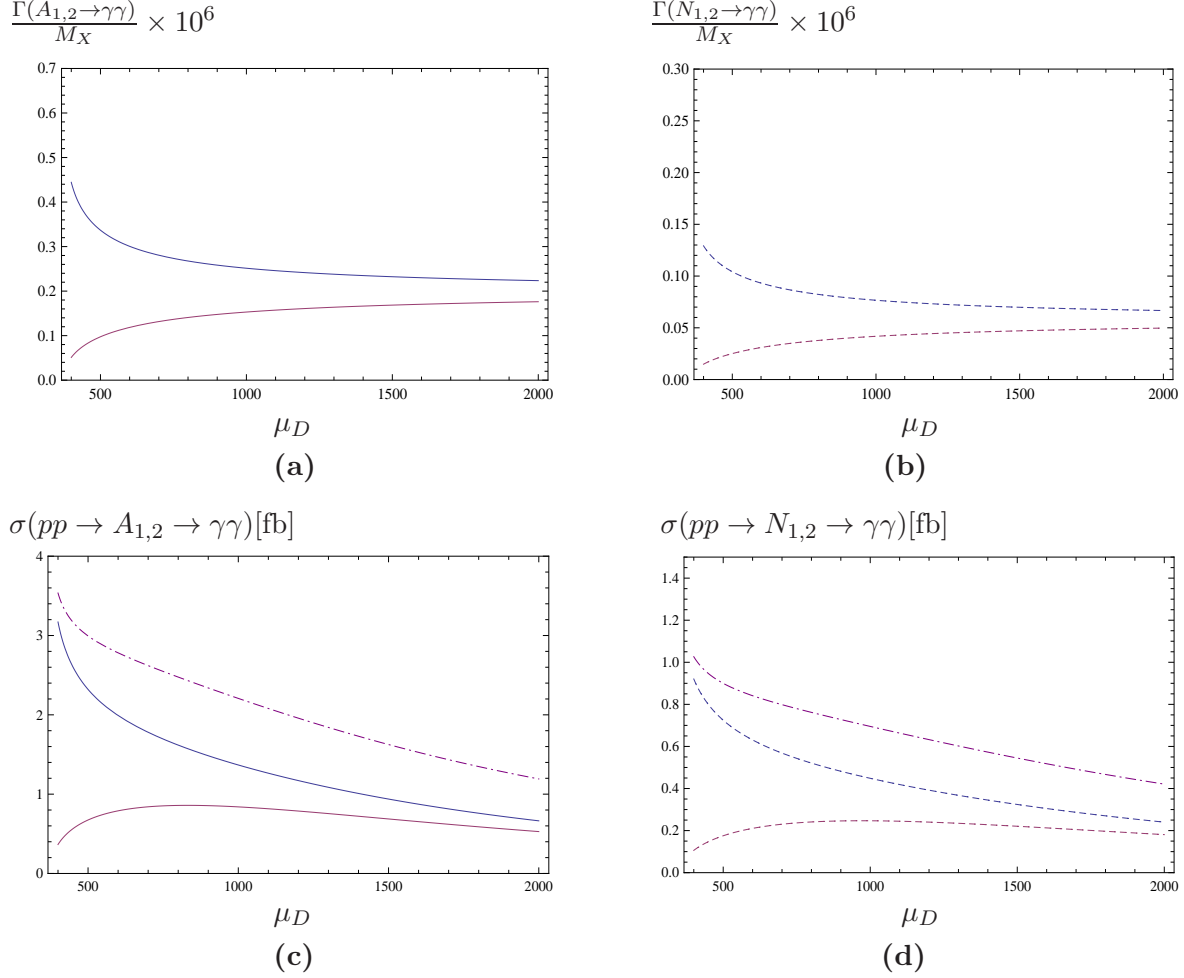


Figure 4: Predictions for the maximal mixing scenario for two maximally mixed pseudoscalars (left panels) or two maximally mixed scalars (right panels) case. In all cases $\mu_{H_\alpha} = 400$ GeV, $\lambda_{\alpha 1} = 0.8$, $\kappa_{i2} = 0.79$, $\lambda_{\alpha 2} = \kappa_{i1} = 0$ and the masses of exotic quarks are set to be equal, i.e. $\mu_{D_i} = \mu_D$. In (a) the ratios $\Gamma(A_1 \rightarrow \gamma\gamma)/M_X$ (upper solid line) and $\Gamma(A_2 \rightarrow \gamma\gamma)/M_X$ (lower solid line) as a function of exotic quark masses in the maximal mixing scenario for $M_{A_\alpha} \simeq 750$ GeV. In (b) the ratios $\Gamma(N_1 \rightarrow \gamma\gamma)/M_X$ (upper dashed line) and $\Gamma(N_2 \rightarrow \gamma\gamma)/M_X$ (lower dashed line) as a function of exotic quark masses in the maximal mixing scenario for $M_{N_\alpha} \simeq 750$ GeV. In (c) the cross sections in fb $\sigma(pp \rightarrow A_1 \rightarrow \gamma\gamma)$ (upper solid line) and $\sigma(pp \rightarrow A_2 \rightarrow \gamma\gamma)$ (lower solid line) as a function of exotic quark masses for $M_{A_{1,2}} \simeq 750$ GeV. The dashed-dotted line correspond to the sum of these cross sections. In (d) the cross sections in fb $\sigma(pp \rightarrow N_1 \rightarrow \gamma\gamma)$ (upper dashed line) and $\sigma(pp \rightarrow N_2 \rightarrow \gamma\gamma)$ (lower dashed line) as a function of exotic quark masses for $M_{N_{1,2}} \simeq 750$ GeV. The dashed-dotted line correspond to the sum of these cross sections.

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